

# Hadron multiplicity in pp and AA collisions at LHC from the Color Glass Condensate

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We provide quantitative predictions for the rapidity, centrality and energy dependencies of inclusive charged-hadron productions for the forthcoming LHC measurements in nucleus-nucleus collisions based on the idea of gluon saturation in the color-glass condensate framework. Our formulation gives very good descriptions of the first data from the LHC for the inclusive charged-hadron production in proton-proton collisions, the deep inelastic scattering at HERA at small Bjorken- $x$ , and the hadron multiplicities in nucleus-nucleus collisions at RHIC.

## I. INTRODUCTION

It has been widely discussed that QCD predicts at high-energy formation of a new state of matter, the so-called color-glass condensate (CGC) [1]. In this state, the density of quarks and gluons  $\rho$  with transverse momenta less than  $Q_s$  reaches a high value, namely  $\rho \propto 1/\alpha_s(Q_s) \gg 1$  where  $\alpha_s$  is the strong coupling constant and  $Q_s$  is a new momentum scale (saturation momentum) that increases with energy. Therefore,  $\alpha_s(Q_s) \ll 1$  and this fact allows us to treat this system on solid theoretical basis. One of the most characteristic and qualitative consequences of the CGC is the emergence of a new mechanism for hadron production at high energy. In this approach, the process of secondary hadron production goes in two stages: production of gluon mini-jet with a typical transverse momentum  $Q_s$ ; and the decay of gluon mini-jets into hadrons.

The CGC picture for the hadron production has passed two critical tests. First, it explains the main features of hadron multiplicity in heavy ion-ion collisions at RHIC (KLN papers [2]); and it predicts the inclusive hadron production in proton-proton ( $pp$ ) collisions [3] at the LHC which was recently confirmed experimentally by the CMS collaboration at  $\sqrt{s} = 7$  TeV [4] and the lower LHC energies [5–7], see also Ref. [8]. The ion-ion collisions at the LHC will lead to a crucial test of the CGC approach. Indeed, the KLN success [2] in the description of RHIC data shows that we are dealing with scatterings of dense partonic systems even at rather low energies. At higher energies we expect that the density of scattering systems increases resulting in more transparent and unambiguous evidence for the creation of the CGC.

In this paper we wish to extend our approach [3] to ion-ion ( $AA$ ) collisions at the LHC. As it was discussed in Ref. [3], we improve the KLN approach [2] in various ways incorporating the unintegrated gluon density that describes HERA data at small Bjorken- $x$ . We also include our knowledge on proton-proton scatterings at the LHC to make our prediction more reliable. In particular, we re-calculate the saturation momentum for nuclei from the saturation momentum in the proton target obtained from DIS data at HERA and confront it with the RHIC's

gold-gold data. On the contrary, in the KLN approach [9], the LHC saturation momentum was found via an extrapolation of the energy dependence of the saturation scale at RHIC in the BFKL region.

In the next section, we discuss the  $k_T$ -factorization and our main formulation for the inclusive hadron production in  $pp$  and  $AA$  collisions. In particular, we introduce the impact-parameter dependent saturation scale for a proton and a nuclear target. Section III is devoted to comparison with the experimental data and to discussion of various predictions for the LHC energies. Finally, we conclude in Sec. IV.

## II. MAIN FORMULATION

The gluon jet production in nucleus-nucleus collisions can be described by  $k_T$ -factorization given by [10],

$$\frac{d\sigma}{dy d^2p_T} = \frac{2\alpha_s}{C_F} \frac{1}{p_T^2} \int d^2\vec{k}_T \phi_A^G(x_1; \vec{k}_T) \phi_A^G(x_2; \vec{p}_T - \vec{k}_T), \quad (1)$$

where  $x_{1,2} = (p_T/\sqrt{s})e^{\pm y}$ ,  $p_T$  and  $y$  are the transverse momentum and rapidity of the produced gluon mini-jet.  $\phi_A^G(x_i; \vec{k}_T)$  denotes the unintegrated gluon density and is the probability to find a gluon that carries  $x_i$  fraction of energy with  $k_T$  transverse momentum in the projectile (or target) nucleus  $A$ . We defined  $C_F = (N_c^2 - 1)/2N_c$  where  $N_c$  denotes the number of colors.

The  $k_T$ -factorization has been proven [10] for the scattering of a diluted system of parton with a dense one. Such a process can be characterized by two hard scales: the transverse momentum of produced particle  $p_T$  and a saturation scale which are both larger than the soft interaction scale  $\mu$ . The ion-ion scattering is a typical example in which we have three scales:  $p_T$  and two saturation scales for the projectile and the target. The  $k_T$ -factorization might be then violated only at low  $x_1$  and  $x_2$ , and when  $p_t$  is smaller than both saturation scales. Outside of this kinematic region we are actually dealing with scatterings of diluted-dense system since one of the saturation scales is small. For the case of scatterings of dense-dense system of partons, the  $k_T$ -factorization has not yet been proven. Nevertheless, the success of the

KLN approach [2] which is based on the  $k_T$ -factorization, in description of the experimental data at RHIC for gold-gold collisions suggests that the  $k_T$ -factorization is currently the best tool that we have at our disposal for the processes of ion-ion scatterings.

The relation between the unintegrated gluon density  $\phi_A^G$  and the color dipole-nucleus forward scattering amplitude has been obtained in Ref. [10]. It reads as follows

$$\phi_A^G(x_i; \vec{k}_T) = \frac{1}{\alpha_s} \frac{C_F}{(2\pi)^3} \int d^2\vec{b} d^2\vec{r}_T e^{i\vec{k}_T \cdot \vec{r}_T} \nabla_T^2 N_A^G(x_i; r_T; b), \quad (2)$$

with notation

$$N_A^G(x_i; r_T; b) = 2N_A(x_i; r_T; b) - N_A^2(x_i; r_T; b), \quad (3)$$

where  $N_A(x_i; r_T; b)$  is the dipole-nucleus forward scattering amplitude which satisfies the perturbative nonlinear small- $x$  Balitsky-Kovchegov (BK) quantum evolution equation [11]. In the above,  $r_T$  denotes the dipole transverse size and  $\vec{b}$  is the impact parameter of the scattering.

Substituting Eq. (2) into Eq. (1), and after analytically performing some integrals, we obtain [3, 10],

$$\frac{d\sigma(y; p_T; \vec{B})}{dy d^2p_T} = \frac{2C_F\alpha_s(p_T)}{(2\pi)^4} \int_{B_1}^{B_2} d^2\vec{B} \int d^2\vec{b} d^2\vec{r}_T e^{i\vec{p}_T \cdot \vec{r}_T} \times \frac{\nabla_T^2 N_A^G(x_1; r_T; b) \nabla_T^2 N_A^G(x_2; r_T; b_-)}{p_T^2 \alpha_s(Q_A(x_1; b)) \alpha_s(Q_A(x_2; b_-))}, \quad (4)$$

where  $\vec{B}$  is the impact parameter between the center of two nuclei,  $\vec{b}$  and  $\vec{b}_- = \vec{b} - \vec{B}$  are the impact parameter between the interacting nucleons with respect to the center of two nuclei. We will see below that we can neglect the impact parameter of the produced gluon-jet from the center of the nucleons.

In the above, we extended the  $k_T$ -factorization given in Eq. (1) by introducing a running strong-coupling  $\alpha_s$ . For the running  $\alpha_s$  we employ the same prescription used in Ref. [3] for  $pp$  collisions. The saturation scale  $Q_A(x_{1,2}, b)$  depends on the  $x_{1,2}$  and the impact parameter and will be introduced in the following. A given centrality bin corresponds to ranges of the impact-parameter  $\vec{B} \in [B_1, B_2]$  of the collisions. The procedure how to select the event with a fixed  $\vec{B}$  and its relation to the centrality bins is well-known, see for example Refs. [9, 12].

Notice that the relation between the unintegrated gluon density and the forward dipole-nucleus amplitude Eqs. (2,3) in the  $k_T$ -factorization Eq. (1) is not a simple Fourier transformation which is commonly used in literature and also depends on the impact-parameter. The impact-parameter dependence in these equations is not trivial and should not be in principle assumed as an overall factor. Using the general properties of high density QCD that the underlying physics depends only on the saturation scale we reconstruct the dipole-nucleus scattering amplitude in two steps. First, we choose the saturation model which effectively incorporates all known saturation properties [3] driven by the BK equation including the impact-parameter dependence of the dipole

amplitude [13]. This model describes both the HERA DIS data at small- $x$  [14] and the proton-proton LHC data [3]. Second, we replace the proton saturation momentum by that of the nucleus. The dipole-nucleus scattering amplitude in our model is given by

$$N_A(x; r; b) = \begin{cases} N_0 \left(\frac{\mathcal{Z}}{2}\right)^{2(\gamma_s + \frac{1}{\kappa\lambda Y} \ln(\frac{2}{\mathcal{Z}}))} & \text{for } \mathcal{Z} \leq 2; \\ 1 - \exp(-\mathcal{A} \ln^2(\mathcal{B}\mathcal{Z})) & \text{for } \mathcal{Z} > 2; \end{cases} \quad (5)$$

where we defined  $\mathcal{Z} = r Q_A(x; b)$ ,  $Y = \ln(1/x)$  and  $\kappa = \chi''(\gamma_s)/\chi'(\gamma_s)$  where  $\chi$  is the LO BFKL characteristic function. The parameters  $\mathcal{A}$  and  $\mathcal{B}$  are determined uniquely from the matching of  $N_A$  and its logarithmic derivatives at  $\mathcal{Z} = 2$ . The nucleus saturation scale is given by

$$Q_A^2(x; b) = \int d^2\vec{b}' T_A(\vec{b} - \vec{b}') Q_p^2(x; b'). \quad (6)$$

In the above, the proton saturation scale  $Q_p$  is defined as

$$Q_p(x; b') = \left(\frac{x_0}{x}\right)^{\frac{\lambda}{2}} \exp\left\{-\frac{b'^2}{4(1-\gamma_{cr})B_{CGC}}\right\}. \quad (7)$$

and  $T_A(B)$  denotes the nuclear thickness. We use for the nuclear thickness the Wood-Saxon parametrization [16]. Notice that for small  $\mathcal{Z} \leq 2$ , the effective anomalous dimension  $\gamma_{cr} = 1 - \gamma_s$  in the exponent in the upper line of Eq. (5) rises from the BK value towards the DGLAP value. The ansatz given in Eq. (5) for a proton target was first introduced in Ref. [15]. The parameters  $\lambda, \gamma_s, N_0, x_0$  and  $B_{CGC}$  are obtained from a fit to the DIS data at low Bjorken- $x$   $x < 0.01$  with a very good  $\chi^2/\text{d.o.f.} = 0.92$  [14].

The saturation scale  $Q_A$  in nuclei is proportional to the density of partons in the transverse plane (see for example Refs. [1, 2, 17]) and Eq. (6) takes this fact into account. In Eq. (6),  $\vec{b}'$  is the impact parameter of the color dipole with respect to the center of the proton,  $\vec{b}$  is the impact parameter of the nucleon with respect to the center of the nucleus and  $\vec{b} - \vec{b}'$  is, therefore, the position of the dipole with respect to the center of the nucleon. The nuclear saturation scale defined in Eq. (6) gives  $Q_A^2 \approx Q_p^2 A^{1/3}$ , in agreement with the basic idea of saturation and the CGC picture [1, 2, 17]. This is because  $T_A \sim A^{1/3}$  (where  $A$  is the effective mass number of the nucleus in a given centrality) and the integral in Eq. (6) is approximately over the nucleon size. Note that since here the nucleon saturation scale depends on impact-parameter, it is not then correct to define the saturation scale on a nuclear target with a simple prefactor  $A^{1/3}$  scaling. In Fig. 1 we show that the impact-parameter dependence of the saturation scale for proton and lead obtained from Eqs. (6,7) are quite different as we expected. Fig. 1 demonstrates the well-known features of nucleus scattering, namely the typical  $b'$  in Eq. (6) is

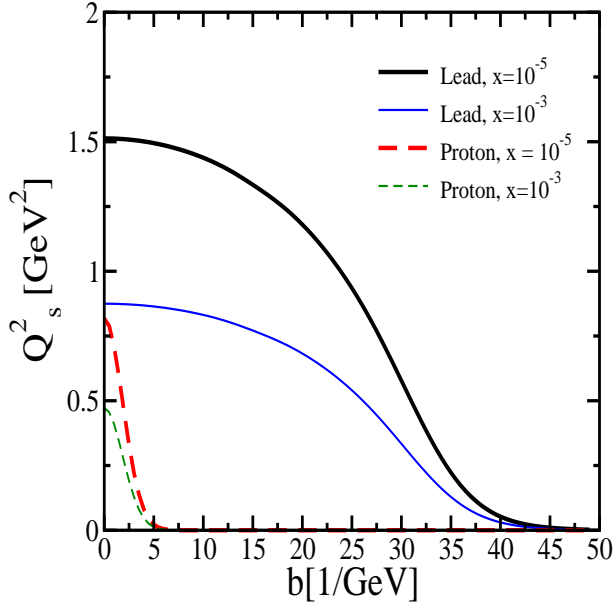


FIG. 1: The impact parameter dependence of the saturation scale for proton and lead at  $x = 10^{-3}$  and  $x = 10^{-5}$ .

much smaller than the typical  $b \sim R_A$  ( $R_A$  is the nuclear radius) in the nuclear thickness.

In order to simulate the behavior of gluon density at large  $x \rightarrow 1$ , we product the unintegrated gluon density with  $(1-x)^4$  as prescribed by quark counting rules and the HERA data on DIS at large- $x$  [3]. Notice that the contribution of  $(1-x)^4$  correction at  $\sqrt{s} = 5.5$  TeV for  $AA$  collisions and  $\eta \leq 4.5$  is less than 4%, and for  $4.5 < \eta < 6$  is less than 13%. Therefore, the main contribution of the unintegrated gluon density at the LHC high-energy in the kinematic region considered here for both  $AA$  and also  $pp$  collisions [3] comes from the small- $x$  region where the saturation physics is important. For the importance of the saturation effects in  $pA$  collisions at the LHC, see Ref. [18].

The  $k_T$ -factorization Eqs. (1,4) gives the cross-section of radiated gluon mini-jets with zero mass while what is actually measured experimentally is the distribution of final hadrons. We therefore should model the non-perturbative hadronization stage of gluon mini-jets. We employ the Local Parton-Hadron Duality principle [19], namely the hadronization is a soft process and cannot change the direction of the emitted radiation. Hence the form of the rapidity distribution of the mini-jet and the produced hadron are different only with a numerical factor  $\mathcal{C}$ . It is well-known that the general assumption about hadronization leads to the appearance of mass of the mini-jet which is approximately on average equal to  $m_{jet}^2 \simeq 2\mu\langle p_T \rangle$  [2, 3] where  $\mu$  is the scale of soft interaction. The mini-jet mass  $m_{jet}$  effectively incorporates the nonperturbative soft pre-hadronization in the pseudo-rapidity space and can be approximately related to the saturation scale [3]. Accordingly, one should also correct the kinematics every where in Eq. (4) due to the

presence of a non-zero mini-jet mass, namely replacing  $p_T \rightarrow \sqrt{p_T^2 + m_{jet}^2}$  in  $x_1, x_2$  and also in the denominator of  $1/p_T^2$ . Finally, in order to take account of the difference between rapidity  $y$  and the measured pseudo-rapidity  $\eta$ , we employ the Jacobian transformation between  $y$  and  $\eta$  [3].

Following Ref. [3], in the spirit of the geometrical-scaling property of the scattering amplitude, we obtain the charged-particle multiplicity distribution at a fixed centrality but various energies from the corresponding mini-jet cross-section Eq. (4) divided by the average area of interaction  $\sigma_s = \mathcal{M}\pi\langle \vec{b}_{jet}^2 \rangle = \mathcal{M}\pi\langle b^2 + b_-^2 \rangle$ . In a similar way, in order to obtain the charged-particle multiplicity distribution at various centrality bins but a fixed energy, we divide the mini-jet cross-section integrated in ranges of  $\vec{B} \in [B_1, B_2]$  with the corresponding relative interaction area  $\sigma_s = \mathcal{M}\pi(B_2^2 - B_1^2)$ .

### III. DISCUSSION AND PREDICTIONS

We have only two unknown parameters in our model: the overall factor  $\frac{KC}{M}$  and the soft-scale  $\mu$  introduced in the definition of mini-jet mas. We also introduced a  $K$ -factor which incorporates the discrepancy between the exact calculation with our formulation. These two phenomenological parameters are fixed at RHIC energy  $\sqrt{s} = 200$  GeV for Au-Au 0 – 6% centrality. Then our results at other energies and various centralities can be considered as free-parameter predictions. The sensitivity of our results to various  $\mu$  is shown in Fig. 2 (top) at  $\sqrt{s} = 200$  GeV. The preferred value of  $\mu$  in Pb-Pb collisions is approximately of order the current-quark mass and is smaller than the corresponding value of  $\mu \approx m_\pi$  in  $pp$  collisions [3]. The decrease in the value of the soft-scale in a denser medium is in accordance with the notion of asymptotic deconfinement that the confinement radius increases with density. The preferred value of the soft-scale in  $AA$  collisions is very close to what was obtained in Ref. [20] for the value of  $\Lambda_{QCD}$  at moderate densities. Loosely-speaking, the smallness of  $\mu$  for  $AA$  collisions compared to the  $pp$  collisions is due to the fact that the mini-jet mass reduces in a denser medium despite the fact that the saturation scale increases. We checked that for Pb-Pb and  $pp$  collisions a fixed mini-jet mass about 0.14 – 0.2 and 0.4 – 0.5 GeV respectively, gives similar results to the case that the mini-jet mass runs with the saturation scale. As we already pointed out the mini-jet mass mimics the properties of the soft nonperturbative pre-hadronization stage and a different value for  $\mu$  in  $AA$  collisions is simply due to the fact that the soft pre-hadronization stage in  $AA$  collisions is different from the corresponding one in  $pp$  collisions.

The main sources of theoretical uncertainties in our approach stem from assuming a fixed  $\mu$  (or a fixed  $m_{jet}$ ) and the prefactor  $\frac{KC}{M}$  for all energies, rapidities and centralities. The experimental errors in the

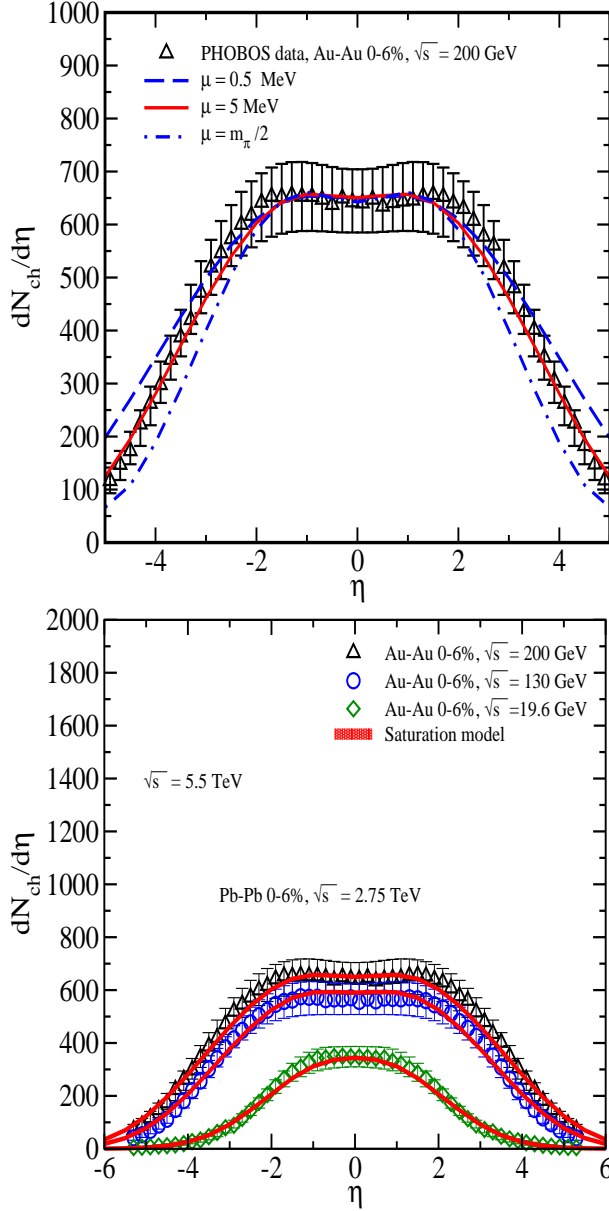


FIG. 2: Top: The effect of the soft-scale  $\mu$  is shown at  $\sqrt{s} = 200$  GeV. Lower: Pseudo-rapidity distribution of charged particles produced in Au-Au and Pb-Pb central 0 – 6% collisions at RHIC  $\sqrt{s} = 19.6, 130, 200$  GeV and the LHC energies  $\sqrt{s} = 2.75, 5.5$  TeV. The band indicates less than 3% theoretical error coming from uncertainties related to normalization and modeling the mini-jet mass. The experimental data are from PHOBOS collaboration [22].

data points taken for fixing these two unknown parameters also induce some uncertainties. Moreover, we employed a Glauber based approach to relate the centrality bins, number of participant and the cut in the impact-parameters [9, 12]. This formulation agreed with the RHIC data on number of participant within error bars, but we expect some uncertainties in the Glauber formulation mainly due to its over-simplicity. In the Glauber

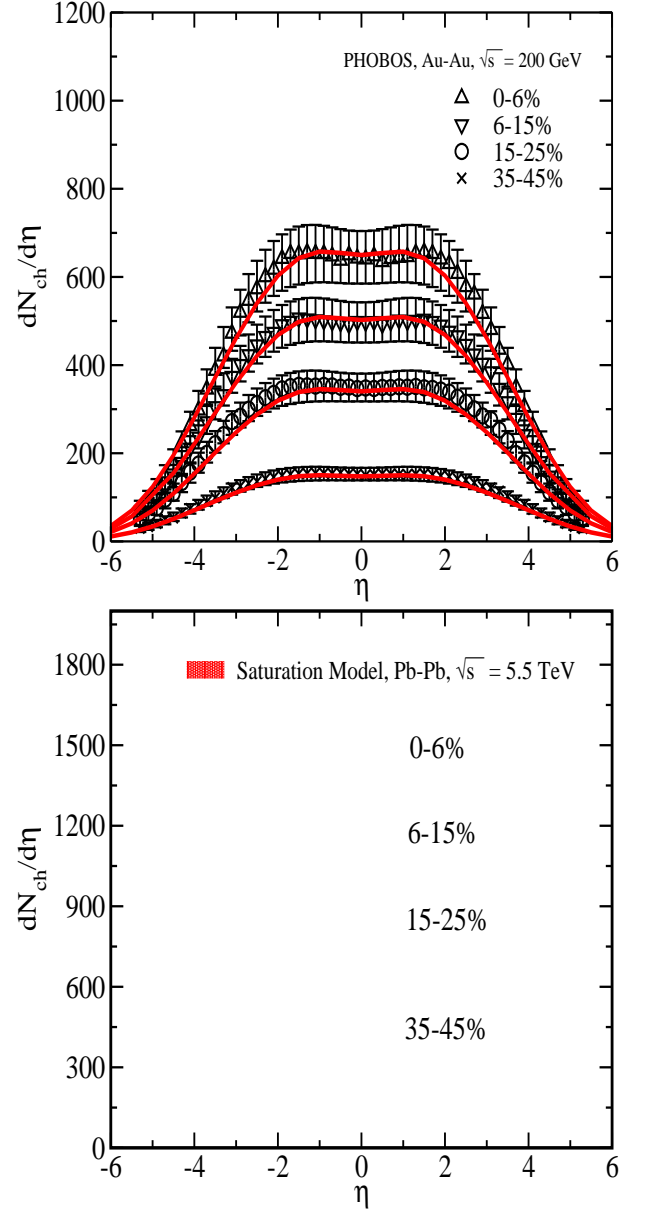


FIG. 3: The pseudo-rapidity dependence at RHIC  $\sqrt{s} = 200$  GeV (top) and the LHC  $\sqrt{s} = 5.5$  TeV (lower) at different centrality bins. The band indicates less than 3% theoretical errors. The experimental data are from PHOBOS collaboration [22].

formalism one should also assume the value of inelastic nucleon-nucleon cross-section  $\sigma_{nn}^{inel}$  (without diffractive components) from outset. Here we used the values obtained in Ref. [21]:  $\sigma_{nn}^{inel} = 64.8, 58.5, 42, 41, 30$  mb for  $\sqrt{s} = 5500, 2750, 200, 130, 19.6$  GeV, respectively. Notice that the nuclear saturation scale defined in Eq. (6) can be in principle different with exact one with extra factor. In order to estimate this uncertainty, we checked that a factor of 2 difference in the definition of Eq. (6) (namely  $Q_A \rightarrow 2Q_A$ ) will change our result less than 5% at the LHC. Overall, including all the above-mentioned possible

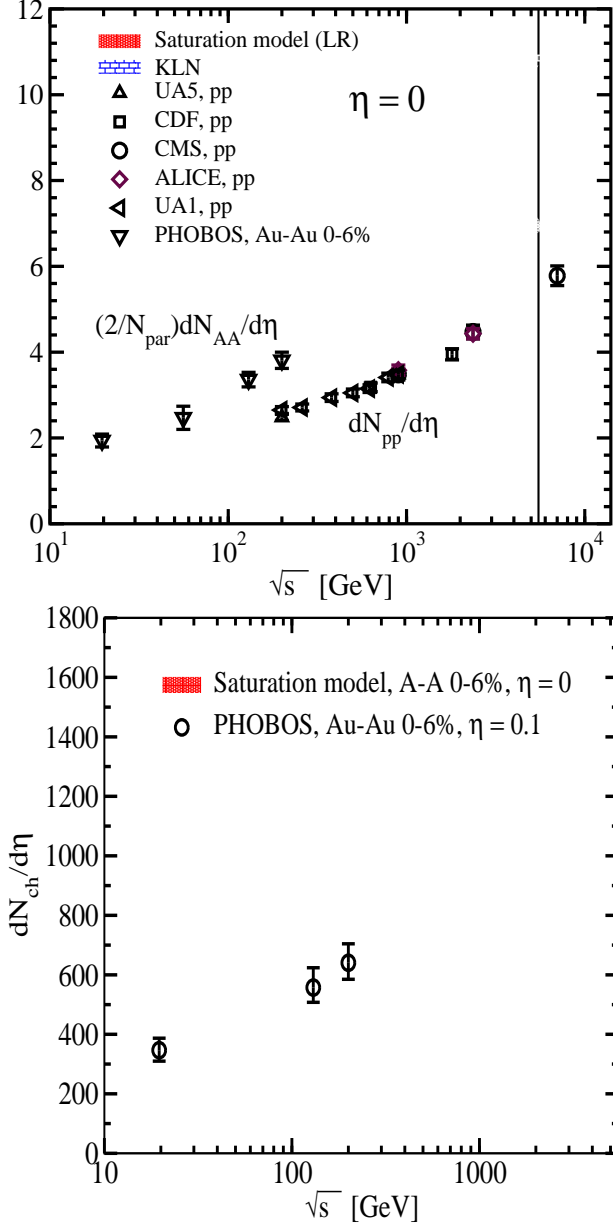


FIG. 4: Energy dependence of the charged hadrons multiplicity at midrapidity  $\eta = 0$  in central collisions in  $pp$  and  $AA$  collisions. The theoretical curve Saturation model (LR) is our prediction. The band indicates less than 3% theoretical errors. The total theoretical uncertainties is less than 7%, see the text for the details. We also show the KLN prediction [9] with the same error band as ours. The experimental data are from [4–6, 22, 23].

theoretical uncertainties, we expect less than 7% theoretical error in our calculations at the LHC energies. The theoretical bar in Fig. 3,4 only show about 3% theoretical error.

In Fig. 2 (lower), we show our predictions at lower RHIC energies  $\sqrt{s} = 19.6$  and 130 GeV in Au-Au collisions, and also for the LHC energies  $\sqrt{s} = 2.75$  and 5.5 TeV in Pb-Pb collisions for 0 – 6% centrality bin.

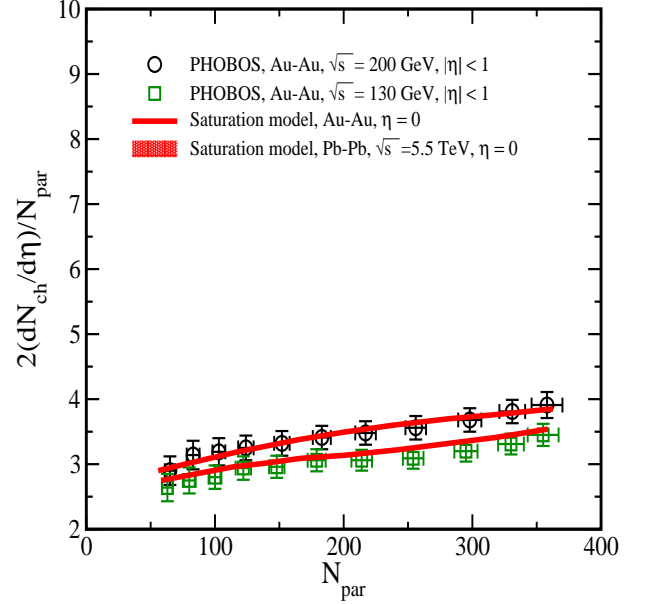


FIG. 5: The scaled pseudo-rapidity density as a function of number of participant  $N_{par}$  at midrapidity for Au-Au at  $\sqrt{s} = 130, 200$  GeV and for Pb-Pb at  $\sqrt{s} = 5.5$  TeV. The experimental data are from PHOBOS collaboration [26].

In Fig. 3, we show the charged-particle multiplicity at various centrality bins for RHIC and the LHC high energy. As we already stressed we only used RHIC data at  $\sqrt{s} = 200$  GeV for 0 – 6% centrality bin in order to fix two unknown parameter of our model. Therefore, the description of other RHIC data here can be considered as predictions. Over-all, our approach gives a very good description of RHIC multiplicity data as a function of rapidity, centrality, number of participant and energy, see Figs. 2-5.

In Fig. 4 we show the energy dependence of  $dN_{pp}/d\eta$ ,  $dN_{AA}/d\eta$  and  $(2/N_{par})dN_{AA}/d\eta$  at midrapidity  $\eta = 0$  for central collisions (where  $N_{par}$  denotes the number of participant for a given centrality). In Fig. 4, we also show our predictions for the charged-hadron multiplicity in  $pp$  collisions. We should emphasize again that the main difference in our formulation for the case of  $pp$  and  $AA$  collisions is only due to the employed different saturation scale for a proton and a nuclear target Eqs. (7,6) (see Fig. 1) and the rest of the formulation is the same. This is in accordance with the notion of universality of the saturation physics which can be further tested at the LHC. In Fig. 4, we show that the recent data from CMS at  $\sqrt{s} = 7$  TeV remarkably confirms our predictions [3] for  $pp$  collisions. Our predictions for  $dN_{AA}/d\eta$  at midrapidity  $\eta = 0$  for central Pb-Pb collisions (for  $B \leq 3.7$  fm or approximately 0 – 6% centrality bin) for  $\sqrt{s} = 2.75$  and 5.5 TeV are  $1152 \pm 81$  and  $1314 \pm 92$ , respectively. Similar results was also suggested in Ref. [24] based on the numerical solution of the BK equation with running strong-coupling with the model assumption that the dipole amplitude does not depend on the impact-parameter. However, in

this paper [24], the relation between the unintegrated gluon-density and the forward dipole-nucleus amplitude in the  $k_T$ -factorization was taken a simple Fourier transformation instead of Eqs. (2,3). The predictions of other approaches at the LHC can be found in Ref. [25].

In Fig. 5, we show  $(2/N_{par})dN_{AA}/d\eta$  as a function of number of participant for RHIC and the LHC energies at midrapidity.

Notice that the KLN approach similar to here is also based on the  $k_T$ -factorization. Both approaches describe the RHIC data but provide rather different predictions at the LHC. In Fig. 4, we also show the KLN predictions for both  $pp$  and  $AA$  collisions. Obviously, the KLN predictions underestimated the  $pp$  multiplicity data at the LHC  $\sqrt{s} = 7$  TeV (in contrast to our predictions) while it overestimates the multiplicity for  $AA$  collisions compared to our predictions. The main differences between our approach and the KLN one is that we used explicitly the impact-parameter dependent form of the  $k_T$ -factorization and employed the correct relation between the unintegrated gluon density and the forward dipole-nucleon amplitude Eqs. (2,3). Then we employed an impact-parameter dependent saturation model which gives a good description of HERA data at small- $x$ . In this sense, we did not have any freedom to model the saturation dipole-proton amplitude at RHIC or the LHC. The key difference between our approach and the KLN one is the fact that we rely on Eq. (6) to determine the saturation scale for nucleus while in the KLN approach [9] the energy dependence was used to obtain the saturation scale for nucleus at the LHC from the one at RHIC.

## IV. CONCLUSIONS

We developed a saturation approach, based on the CGC theory, which describes DIS data at HERA [14], hadron production at the LHC [3] in  $pp$  collisions and at RHIC in  $AA$  collisions (this paper). In this approach we treat in the same way proton and nuclear target considering that the only difference between them is the value of the saturation scale. The nonperturbative stage of the jet-hadronization was described in the same scheme for both proton and nuclear target, minimizing the uncertainties in prediction that could stem from this stage. We provided here various quantitative predictions including the rapidity, centrality and energy dependencies of the inclusive charged-hadron multiplicity in  $AA$  collisions at the LHC. We believe that our predictions for nucleus-nucleus collisions at the LHC will be a crucial test of the CGC approach.

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